

KEK Preprint 2010-41

FTUV-10-1129

October 25, 2011

Measuring Anomalous Couplings in $H \rightarrow WW^*$ Decays at the International Linear Collider

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Abstract

Measurement of the Higgs coupling to W -bosons is an important test of our understanding of the electroweak symmetry breaking mechanism. We study the sensitivity of the International Linear Collider (ILC) to the presence of anomalous HW^+W^- couplings using $ZH \rightarrow \nu\bar{\nu}WW^* \rightarrow \nu\bar{\nu}4j$ events. Using an effective Lagrangian approach, we calculate the differential decay rates of the Higgs boson including the effects of new dimension-5 operators. We present a Monte Carlo simulation of events at the ILC, using a full detector simulation based on geant4 and a real event reconstruction chain. Expected constraints on the anomalous couplings are given.

1 Introduction

With the recent announcement of high-energy collisions at the Large Hadron Collider (LHC), discovery of the long-awaited Higgs boson is eagerly anticipated. The Higgs field is crucial to the Standard Model (SM), as it provides the mechanism for electroweak (EW) symmetry breaking and gives rise to the masses of both the W^\pm and Z gauge bosons and all charged fermions when it develops a vacuum expectation value (VEV).

In practice, the discovery of a SM-like Higgs boson will be only the beginning of the story. Over the last thirty years, a number of alternative models have been proposed, notably low-energy supersymmetry and models with extra dimensions, which attempt to resolve the so-called hierarchy problem of the SM. Once discovered, it will be crucial to measure the properties of the Higgs particle in order to ascertain to which, if any, of the proposed models it belongs. Of particular importance will be the task of verifying that the Higgs boson we have discovered really is the particle responsible for EW symmetry breaking and the generation of fermion masses.

Indeed, it is notable that one of the “gold plated” Higgs boson discovery channels at the LHC involves production of the Higgs particle through gluon fusion, followed by its decay into photon pairs, i.e. $gg \rightarrow H \rightarrow \gamma\gamma$. This channel involves loop-induced couplings in both production and decay and the discovery of such a resonance tells us very little about EW symmetry breaking!¹ For this, accurate measurements of the Higgs boson couplings to the EW gauge bosons must be made. The LHC can be used to extract some information on the Higgs boson coupling to Z bosons using the decay $H \rightarrow ZZ \rightarrow (l^+l^-)(l'^+l'^-)$ [3], since this final state can be efficiently triggered; effects of anomalous HW^+W^- couplings can also be probed through their contribution to WW scattering [4] and the gauge-boson fusion Higgs production mechanism [5]. For a direct measurement of the HW^+W^- coupling, however, the best environment is a lepton collider, such as the International Linear Collider (ILC).

In an electron-positron collider such as the ILC, Higgs bosons are predominantly produced through EW interactions, either through Higgs-strahlung from virtual Z -bosons or through gauge boson fusion; such reactions can therefore be used to probe anomalous Higgs-gauge-gauge couplings [6]. Furthermore, the clean environment of

¹ In fact, a scalar in a certain class of new physics models [1], which is nothing to do with the EW symmetry breaking, can mimic such a signal. It has been studied [2] how well the International Linear Collider distinguishes the scalar from the Higgs boson.

a lepton collider also allows for studies based on the asymmetries in the decays of the Higgs boson [7]. For a heavy Higgs boson (with mass $\gtrsim 200$ GeV), the authors of [8] have proposed to accurately measure the couplings instead at a future photon collider.

If, as we expect, the underlying theory is gauge invariant, then anomalous HWW couplings necessarily imply anomalous contributions to the HZZ , $H\gamma\gamma$ and mixed $HZ\gamma$ vertices. Whilst measurements of anomalous Higgs boson couplings to the neutral vector bosons may make use of the very high rates of the Higgs-strahlung process, determining the structure of the HW^+W^- vertex relies on either the gauge-boson fusion production process or on the decay $H \rightarrow WW^*$. Furthermore, these measurements should be performed independently of one another, as a test of the underlying gauge invariance. In this paper we concentrate on the decay process for two reasons; firstly, the differential cross sections considered do not depend on the additional HNN , $N = Z, \gamma$ anomalous couplings at leading order and secondly, this process allows us to study the effects of CP-violating parameters. In contrast, since the final state forward neutrinos are not measured in the gauge-boson fusion process $e^+e^- \rightarrow H\nu\bar{\nu}$, this process has no sensitivity to CP-violating parameters; furthermore, there are large backgrounds from the related process $e^+e^- \rightarrow He^+e^-$, where the final state electrons also escape detection, which introduces a significant dependence on the HNN anomalous couplings.

In particular, we study the feasibility of measurements for anomalous Higgs boson couplings to W^+W^- pairs using ZH production followed by $H \rightarrow WW^*$ and $Z \rightarrow \nu\nu$ at the ILC, based on realistic Monte Carlo simulations. We stress that this is the first full simulation study of its kind using a full detector simulator based on geant4 and a real event reconstruction chain. The Higgs production mechanism clearly depends on the presence of anomalous HZZ and $HZ\gamma$ couplings, however the distributions of the Higgs boson decay products which constitute our signal do not. The only possible effect of anomalous HNN couplings in the production mechanism is a change in the overall rate $e^+e^- \rightarrow ZH \rightarrow ZWW^*$ and can easily be measured at relatively low integrated luminosities [6]. Hence, we feel justified in neglecting the effect of these couplings in our analysis focusing on the Higgs decay $H \rightarrow WW^*$.

The structure of our paper is as follows. In the next section we outline the effective interaction Lagrangian under discussion and present analytic formulas for the relevant differential decay rates. In Section 3, we present details of the Monte Carlo simulation and results of the analysis. We discuss these results and a simple

example of new physics model which give rise to the effective interaction Lagrangian in Section 4. The final section is reserved for summary and conclusions.

2 Physics Model

We may parametrise the relevant terms of the general interaction Lagrangian, which couples the Higgs boson to EW vector bosons in a both Lorentz- and gauge-symmetric fashion, as

$$\mathcal{L}_{\text{HWW}} = 2M_W^2 \left(\frac{1}{v} + \frac{a}{\Lambda} \right) H W_\mu^+ W^{-\mu} + \frac{b}{\Lambda} H W_{\mu\nu}^+ W^{-\mu\nu} + \frac{\tilde{b}}{\Lambda} H \epsilon^{\mu\nu\sigma\tau} W_{\mu\nu}^+ W_{\sigma\tau}^- , \quad (2.1)$$

where M_W is the mass of the W -boson, $W_{\mu\nu}^\pm$ is the usual gauge field strength tensor, $\epsilon^{\mu\nu\sigma\tau}$ is the Levi-Civita tensor, v is the VEV of the Higgs field, a, b, \tilde{b} are real dimensionless coefficients and Λ is a cutoff scale. The SM interaction is recovered in the limit $a, b, \tilde{b} \rightarrow 0$. The dimensionless couplings b, \tilde{b} parametrise the leading dimension-five non-renormalisable interactions², which we assume are due to contributions arising from some new physics at the scale Λ . The dimensionless coupling a represents corrections to the SM term, assumed to originate at the same scale Λ . The Lagrangian (2.1) is not by itself gauge invariant; to restore explicit gauge invariance we must also include the corresponding anomalous couplings of the Higgs boson to Z bosons and photons.

We will assume the Higgs boson mass to be $M_H < 2M_W$ so that the decay to real W^+W^- pairs is kinematically forbidden; the anomalous couplings may however contribute to the decay $H \rightarrow WW^*$ with distinct signatures. The a parameter is simply a rescaling of the SM coupling and therefore manifests itself as a shift in the overall partial width for this channel. By comparison, the non-renormalisable coupling b has a different Lorentz structure to the SM term and leads to a change in the ratio of couplings to the transverse or longitudinal components of the gauge bosons. Finally, the coupling \tilde{b} introduces a CP-violating operator which can affect angular correlations, as discussed below.

Assuming all final state fermions to be massless, the differential partial width for the decay chain $H \rightarrow WW^* \rightarrow 4j$ as a function of the on-shell W -boson momentum p_W and the azimuthal angle between the up-type quark and anti-quark ϕ (with axis

²The effects of dimension-six operators in the effective Lagrangian were considered in [9].

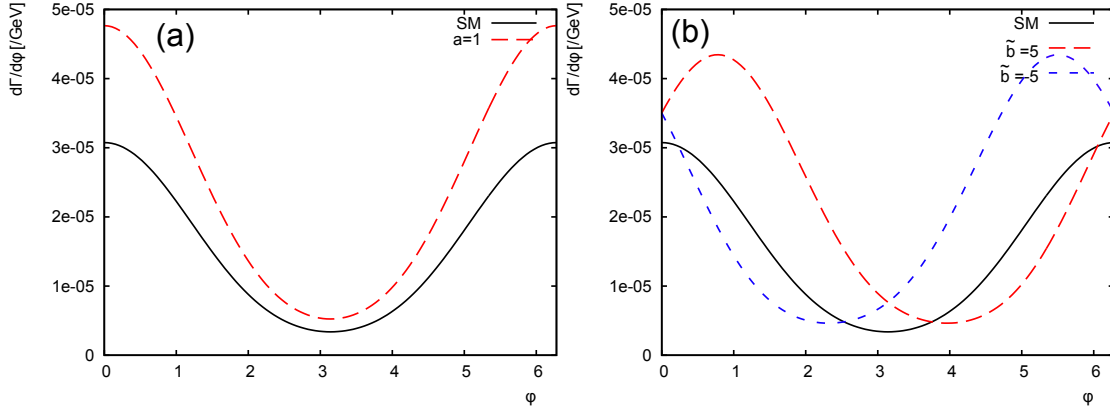


Figure 1: The ϕ -distribution of the decay $H \rightarrow WW^* \rightarrow 4j$ with the inclusion of anomalous couplings. (a) The SM curve along with that for $a = 1$, $b = \tilde{b} = 0$, same for both distributions. (b) The SM result with the cases $\tilde{b} = \pm 5$, $a = \tilde{b} = 0$, $\Lambda = 1$ TeV; the position of the minimum is now shifted as discussed in the text.

of rotation in the direction of the W^+ momentum) is given by

$$\frac{d^2\Gamma}{dp_W d\phi} = \left(N_c \sum_{\substack{i=u,c \\ j=d,s,b}} |V_{ij}|^2 \right)^2 \frac{1}{2M_H} \frac{1}{2M_W \Gamma_W} \frac{M_W^{10} k^2 p_W^2}{(k^2 - M_W^2)^2 \pi^5 E_W v^6} \times \left[\frac{c_0}{36} + \left(\frac{c_\phi e^{i\phi} + c_\phi^* e^{-i\phi}}{512} \right) + \left(\frac{c_{2\phi} e^{i2\phi} + c_{2\phi}^* e^{-i2\phi}}{288} \right) \right], \quad (2.2)$$

where $N_c = 3$ is the number of colors, V_{ij} is the Cabibbo-Kobayashi-Maskawa quark mixing matrix, Γ_W is the W -boson width, $E_W = \sqrt{M_W^2 + p_W^2}$ is the energy of the on-shell W -boson and $k^2 = (M_H - E_W)^2 - p_W^2$ is the invariant squared mass of the off-shell W -boson. The coefficient functions c_0 , c_ϕ , $c_{2\phi}$ can be written in terms of two dimensionless combinations of parameters,

$$\begin{aligned} c_0 &= |T|^2 + \frac{1}{2} L^2, \\ c_\phi &= \pi^2 T L, \\ c_{2\phi} &= T^2, \end{aligned} \quad (2.3)$$

where the real function L is due to contribution from longitudinally polarized on-shell W -bosons and the complex function T from transversely polarized bosons.

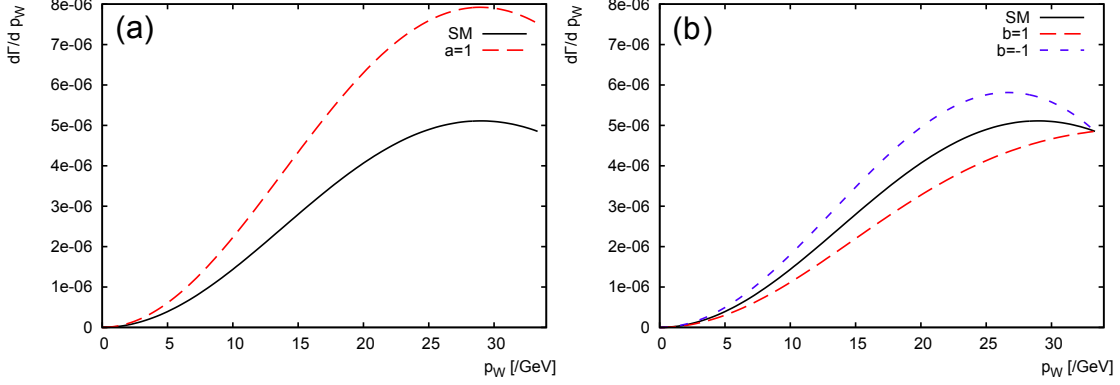


Figure 2: The p_W distribution of the decay $H \rightarrow WW^* \rightarrow 4j$ with the inclusion of anomalous couplings. (a) The SM curve along with that for $a = 1$, $b = \tilde{b} = 0$, $\Lambda = 1$ TeV. (b) The SM result with the cases $b = \pm 1$, $a = \tilde{b} = 0$, $\Lambda = 1$ TeV.

Explicitly, they are given by

$$\begin{aligned}
T &= \left(1 + \frac{av}{\Lambda}\right) - \frac{bv}{\Lambda} \frac{(p_W^\mu k_\mu)}{M_W^2} - i \frac{2M_H p_W}{M_W^2} \frac{v \tilde{b}}{\Lambda}, \\
L &= \left(1 + \frac{av}{\Lambda}\right) \frac{(p_W^\mu k_\mu)}{M_W k} - \frac{bv}{\Lambda} \frac{\left((p_W^\mu k_\mu)^2 - M_H^2 p_W^2\right)}{M_W^3 k},
\end{aligned} \tag{2.4}$$

where k_μ is the 4-momentum of the off-shell W -boson and $k = \sqrt{k_\mu k^\mu}$. The 4-vector product $p_W^\mu k_\mu$ can be expanded as $p_W^\mu k_\mu = M_H E_W - M_W^2$.

We see that for $\tilde{b} = 0$, all coefficients c_0 , c_ϕ , $c_{2\phi}$ are real and the partial decay width is a function of cosines only. In the SM limit, the magnitudes of the transverse function T and the longitudinal function L are approximately equal, assuming a W -boson energy of the order of the W mass, $E_W \sim M_W$. From (2.3), in this limit we have the ratio of coefficients $c_\phi/c_{2\phi} \sim \pi^2$ so that the $\cos \phi$ term dominates, the minimum of the distribution is seen to be at $\phi_{\min} = \pi$. Non-zero values of \tilde{b} shift the minimum of the ϕ distribution to $\phi_{\min} = \pi + \delta$ with $\delta = \arg(T)$. This effect is illustrated in Fig. 1, where we plot the ϕ -dependence of the partial width in both the SM and taking $a = 1, b = \tilde{b} = 0$ (Fig. 1 (a)) and $a = 0, \tilde{b} = \pm 5$ (Fig. 1 (b)), with $\Lambda = 1$ TeV.

The energy of the on-shell W boson in the decay $H \rightarrow WW^*$ is not fixed by kinematic constraints. After integrating (2.2) over ϕ we see that only the c_0 term contributes to the differential decay rate $d\Gamma/dp_W$. The presence of the anomalous couplings b, \tilde{b} modifies the energy-dependence of this expression through the L and

T functions, as shown in Fig. 2, which compares the effect of non-zero a and b terms. In particular, the contribution of b is seen to vanish at the kinematic limit of the distribution.

3 Monte Carlo Simulation

3.1 Simulation Conditions and Tools

For the simulation of the measurement of the anomalous HWW couplings, we take the Higgs mass to be 120 GeV along with a center of mass energy of 250 GeV and an integrated luminosity of 250 fb^{-1} as assumed in the Letters of Intent (LoI) for the ILC detectors [10, 11, 12]. Notice that for a Higgs mass of 120 GeV the ZH cross-section attains its maximum at around this energy³ and the branching ratio for $H \rightarrow WW^*$ decay is 15.0%, which is subdominant. The angular analyses and W momentum measurement discussed in the previous section necessitate the identification of the four jets from the $H \rightarrow WW^*$ decay and their correct pairing. Given the branching fraction of this decay mode and the additional combinatorial background in the jet paring which would otherwise hamper the analyses, we require the associated Z to decay into $\nu\bar{\nu}$. Our signal thus consists of four jets plus missing energy. Consequently, four fermion final states such as $\nu\nu qq$, $qql\nu$, $\ell\ell\ell\ell$, $qq\ell\ell$, and $qqqq$, primarily coming from WW and ZZ production, comprise the SM background. In order to suppress these backgrounds, which would otherwise dominate in this study, we use 80% right-handed polarization for the electron beam and 30% left-handed polarization for the positron beam in the following analysis.

The signal events $ZH \rightarrow \nu\nu WW^*$ were generated using *Physsim* [13], with the cutoff scale Λ taken to be 1 TeV when the anomalous couplings were switched on⁴. $ZH \rightarrow \nu\nu H$ events for decay modes other than $H \rightarrow WW^*$ were generated by *WHIZARD*, together with all the other SM backgrounds. In both generators, initial-

³ Note also that the mass resolution for the Higgs boson recoiling against lepton pairs from the accompanying Z boson is known to degrade with the energy. The most accurate determination of the Higgs mass is hence expected in this energy region, by investing a substantial running time to accumulate an integrated luminosity of 250 fb^{-1} , together with a model-independent determination of the total ZH production cross section, which is indispensable to measurements of the various branching ratios of the Higgs boson.

⁴ Note that the absolute values of the anomalous couplings such as a , b , and \tilde{b} become meaningful only after the cutoff scale is given.

state radiation and beamstrahlung have been included in the event generation. The beam energy spread was set to 0.28% for the electron beam and 0.18% for the positron beam. We have ignored the finite crossing angle between the electron and positron beams. In the event generation, helicity amplitudes were calculated using the HELAS library [14], which allows us to deal with the effect of gauge boson polarizations properly. The event generator Pythia6.409 [15] was used for parton-showering and hadronization. The generator data for the SM background had been prepared as a common data sample for the LoI studies and stored in the StdHep format [16] at SLAC. The SM background sample consists of all the SM processes with up to 4 fermions in the final state, which is about 10 million events in total. Since the cross-section of 6 fermion events is small compared to that of the signal at $\sqrt{s} = 250$ GeV and can be rejected easily, we ignored these events.

The 4-momenta of the (quasi-)stable particles after parton-showering and hadronization were fed into a geant4-based full detector simulator called Mokka [17], in which ILD_00 is implemented as the detector model [10]. The ILD_00 detector model consists, from inside to outside, of a very thin 6-layer vertex detector with a point resolution of about $3\mu\text{m}$, silicon internal and forward trackers, a time projection chamber (TPC) having about 200 sample points with a point resolution of $100\mu\text{m}$ or better, silicon external and endcap trackers, ultra high granularity electro-magnetic and hadron calorimeters, a superconducting solenoid of 3.5 T, and return yokes interleaved with muon detectors. With this detector model, we expect the transverse momentum resolution ($\Delta(1/p_T) = \Delta p_T/p_T^2$) to be $2.0 \times 10^{-5} \text{ GeV}^{-1}$ asymptotically, rising to $9.0 \times 10^{-5} \text{ GeV}^{-1}$ at 10 GeV, and to $9.0 \times 10^{-4} \text{ GeV}^{-1}$ at 1 GeV.

The generated detector hits and signals were processed through a real event reconstruction program called MarlinReco implemented in the Marlin framework [18]. In the event reconstruction, charged particle tracks were reconstructed from tracker hits by a realistic track finder and a Kalman-filter-based track fitter, taking into account signal overlapping as well as energy loss and multiple scattering. Calorimeter hits were then clustered and combined with the tracker information to perform a particle flow analysis (PFA) [19] to achieve the best jet energy resolution. The jet energy resolution for 45 GeV jets from $Z \rightarrow q\bar{q}$ events is estimated to be 3.7%, which improves to about 3% for 100 GeV jets. Jet clustering was done with the Durham algorithm [20], and the resultant jets were flavor-tagged on a jet-by-jet basis with the LCFIVertex package [21] after vertex finding with the ZVTOP algorithm [22].

3.2 Event Selection

The goal of our event selection is to isolate the signal events with four jets plus missing energy originating from ZH production followed by $H \rightarrow WW^*$ and $Z \rightarrow \nu\bar{\nu}$ decays. We thus started our event selection by forcing all the events to cluster into four jets by adjusting the Y_{cut} value [20]. The Higgs boson and on-shell W boson masses were then reconstructed by paring these four jets so as to minimize the χ^2 function defined by

$$\chi^2 = \frac{(\text{rec} M_H - M_H)^2}{\sigma_H^2} + \frac{(\text{rec} M_W - M_W)^2}{\sigma_W^2}, \quad (3.5)$$

where $\text{rec} M_H$ is the reconstructed Higgs mass, M_H is the input Higgs mass (120 GeV), $\text{rec} M_W$ is the reconstructed on-shell W mass, M_W is the nominal W mass (80.4 GeV) and $\sigma_{H(W)}$ is the mass resolution for the Higgs(W).

After the mass reconstruction, we required the reconstructed Higgs mass ($\text{rec} M_H$) to lie in the range $110 \text{ GeV} < \text{rec} M_H < 130 \text{ GeV}$. Since we assume a Z boson decaying into a neutrino pair in the signal event, resulting in a missing mass peak at the Z mass, we required a missing mass in the range $70 \text{ GeV} < \text{miss} M < 140 \text{ GeV}$.

The main backgrounds in this analysis are from $e^+e^- \rightarrow WW$ and ZZ . The angular distributions of these processes have peaks in the forward and backward regions. For this reason, we required the angle of the reconstructed Higgs boson with respect to the beam axis ($\cos \theta_H$) to be $|\cos \theta_H| < 0.95$. We then looked at the Y -value for the forced 4-jet clustering, which is expected to be small for $\nu\nu qq$ and $\nu\nu \ell\ell$ events having only two ‘‘partons’’ in their final states. We therefore selected events with $Y_- > 0.0005$, where Y_- is the threshold Y -value at which the number of jets changes four to three. After the selection cuts described so far, the dominant background became $\ell\nu qq$. The lepton in the $\ell\nu qq$ final state comes from the leptonic decay of a W and has a larger energy than leptons from jets. We hence required the maximum track energy (E_{trk}) to be below 30 GeV.

Since we are focusing our attention on the $H \rightarrow WW^*$ decay in this analysis, the $ZH \rightarrow \nu\nu bb$ process is a background to be discarded. We thus rejected $ZH \rightarrow \nu\nu bb$ events by requiring the number of b -tagged jets ($^{4\text{-jet}} N_b$) to be $^{4\text{-jet}} N_b \leq 1$. Since the $ZH \rightarrow \nu\nu bb$ channel has two jets in the final state, candidate events were further jet-clustered into two jets and then required to have no b -tagged jets, $^{2\text{-jet}} N_b = 0$.

After all the selection cuts, we performed a likelihood analysis as follows. We used $\text{miss} M$, $\cos \theta_H$, Y_- , $^{4\text{-jet}} N_b$, and the number of charged tracks as the input variables

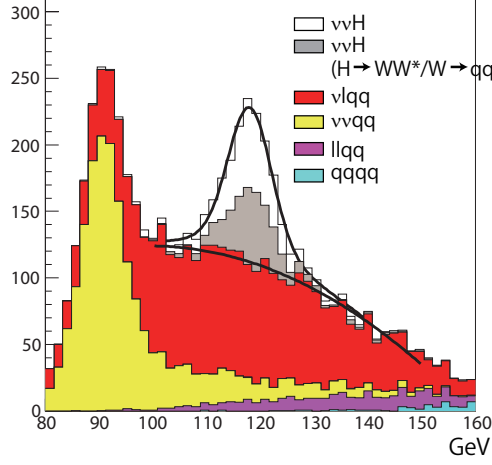


Figure 3: Distribution of the reconstructed Higgs mass after selection cuts.

Process	No cut	After cuts	$\mathcal{L}_{\text{cut}} > 0.79$	$N_c = 2$
$\nu\nu H(H \rightarrow \text{all})$	10,634	1,518	756	546
$\nu\nu H(H \rightarrow WW^* \rightarrow 4\text{-jet})$	680	512	348	258
$llll$	753,964	46	0	0
$qqqq$	378,726	8	3	2
$llqq$	335,762	409	94	70
νlqq	299,866	8,571	1,063	692
$\nu\nu ll$	103,704	3	0	0
$\nu\nu qq$	63,649	1,090	207	110

Table 1: Cut summary.

of the likelihood function and tuned the likelihood cut position to maximize signal significance. We obtained a maximum signal significance of 7.6 at a likelihood cut position $\mathcal{L}_{\text{cut}} = 0.79$. Figure 3 shows the reconstructed Higgs mass distribution after all cuts. Fitting the distribution with a double Gaussian plus a second order polynomial, we estimated the expected accuracy of the branching ratio $\text{BR}(H \rightarrow WW^*)$ to be 15.7%, assuming that the measurement accuracy of the ZH cross-section is 2.5% [10]. The branching ratio can be determined to an accuracy of 5%, however, by using processes with leptonic decays of the W [23]. The number of events before and after the selection cuts are summarized in Table 1.

3.3 Analysis Results

We investigated the distributions of the variables sensitive to the anomalous couplings. The distributions of the W boson momenta in the Higgs rest-frame (p_W) and the jet angle in the W boson rest-frame were plotted after selection cuts. The jet angle distributions were plotted for the on-shell ($\cos\theta_{j1}$) and off-shell W bosons ($\cos\theta_{j2}$) separately. We also examined the distribution of the angle between the two W boson decay planes (ϕ_{plane}) corresponding to that between the two up-type quarks from the decays of the W bosons. We applied double c -tagging to select the two up-type quarks (c quark) in $ZH \rightarrow \nu\nu WW^* \rightarrow \nu\nu cscs$, where the selection efficiency was 88% as shown in Table 1. The ϕ_{plane} was histogrammed using the two c -tagged jets, without identifying their charges. Distributions for the ZH events were obtained, evaluating the contamination from the SM backgrounds by fitting the Higgs mass distribution for each ϕ_{plane} bin. Since the branching ratio of the Higgs to channels other than $H \rightarrow WW^*$ will be determined with much better accuracies than the statistical errors shown on the distributions [23], we subtracted the background from these decay modes ignoring their systematic errors on the cross-section to obtain the distributions of p_W , $\cos\theta_{j1}$, $\cos\theta_{j2}$, and ϕ_{plane} for $ZH \rightarrow \nu\nu WW^*$ events.

To estimate the sensitivity to the Higgs anomalous couplings, the distributions of p_W , $\cos\theta_{j1}$, $\cos\theta_{j2}$, and ϕ_{plane} for events with non-zero anomalous couplings were compared with the SM case. For the comparison we varied two of the parameters a , b , and \tilde{b} , whilst the third was set to zero. We then drew probability contours for $\Delta\chi^2 = 1$, 2.28 and 5.99, corresponding to 39%, 68% and 95% confidence levels (C.L.) respectively, as in Figs. 4 through 6.

The contour plot in the a - b plane (Fig. 4) shows a linear correlation between a and b due to changes in absolute value of the $ZH \rightarrow \nu\nu WW^*$ cross section, which increases with a but decreases with increasing b . Note that with our conventions $a \simeq -4.1$ cancels the SM coupling of the Higgs to W^+W^- , which means that taking $a \simeq -8.2$ effectively reverses the sign of the SM coupling term. If we reverse the sign of the b term in addition, we hence obtain exactly the same distribution provided that the other parameter, \tilde{b} , is kept at zero. For this reason we observe a second allowed region in Fig. 4, connected to the first region (containing the SM point $a = b = 0$) by a 180° rotation about $(a, b) = (-4.1, 0)$.

By the same token, we have two allowed regions for the contours in the a - \tilde{b} plane

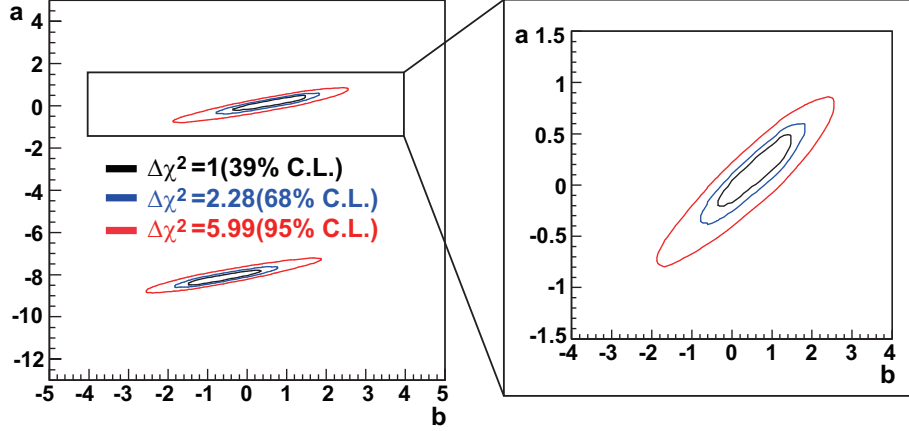


Figure 4: Probability contours for $\Delta\chi^2 = 1, 2.28$, and 5.99 in the a - b plane, which correspond to 39%, 68%, and 95% C.L., respectively.

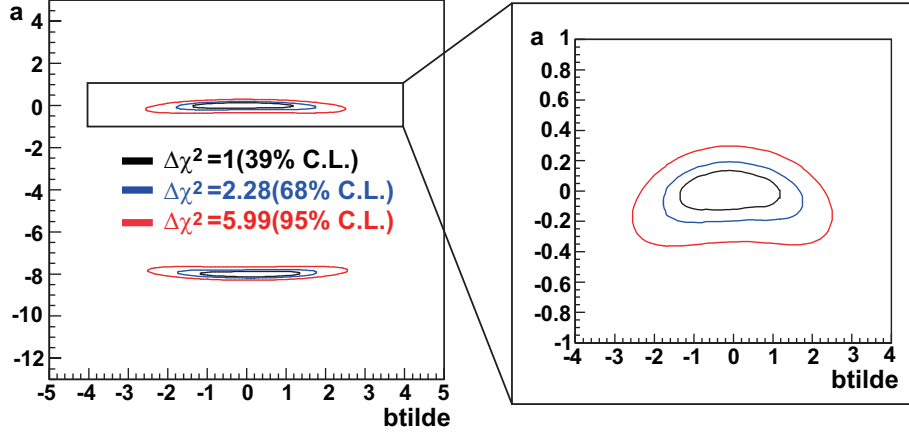


Figure 5: Contours similar to Fig. 4 plotted in the a - \tilde{b} plane.

as plotted in Fig. 5. The additional mirror symmetry of the contours about $\tilde{b} = 0$ is present because we did not identify the charge of the charm jets for the ϕ_{plane} measurement. The prospects for resolving this additional degeneracy by measuring the jet charge are discussed in Section 4.1.

Finally, Fig. 6 shows the contours in the b - \tilde{b} plane for $a = 0$. We observe that these contours are also symmetric under the replacement $\tilde{b} \rightarrow -\tilde{b}$, again due to the non-identification of the jet charge.

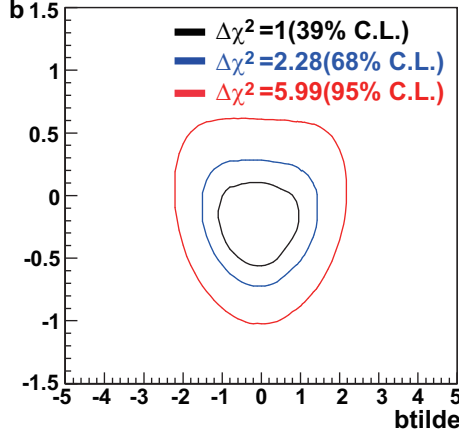


Figure 6: Contours similar to Fig. 5 plotted in the b - \tilde{b} plane.

4 Discussions

4.1 Charge identification of c -jets

If the charge of c -jets can be identified, the shape of the ϕ_{plane} distribution will change depending on the sign of \tilde{b} . For this reason, it might be possible to identify the sign of \tilde{b} by using the identification of c -jets. Since the charge identification of b - and c -jets is considered possible at the ILC by measuring the number of positive and negative charged tracks in the jet clusters, we investigated the sensitivity to the anomalous coupling including identification of the jet charge.

The efficiency of the charge identification for c -jets was assumed to be 14.6% [21] and the same efficiency was used for the backgrounds. We ignored the probability to mis-identify the jet charge as the opposite sign. If we require that the charge of at least one c -jet candidate is identified, 27% of $ZH \rightarrow \nu\nu WW^* \rightarrow \nu\nu cscs$ events can be selected. By using these events, the ϕ_{plane} distribution was produced considering the relative direction of the positive and negative charge (ϕ_{plane}^{+-}). If the charge of both c -jet candidates was not identified, ϕ_{plane} was calculated just as the angle between two c -jet candidates (ϕ_{plane}^{00}). Then, the sensitivity to the anomalous coupling was evaluated by using p_W , $\cos\theta_{j1}$, $\cos\theta_{j2}$, ϕ_{plane}^{+-} , and ϕ_{plane}^{00} .

Figure 7 shows the probability contours for $\Delta\chi^2 = 1, 2.28$, and 5.99 in the a - \tilde{b} plane. Here, the integrated luminosity is taken to be $1,000 \text{ fb}^{-1}$, since the statistics of $ZH \rightarrow \nu\nu WW^* \rightarrow \nu\nu cscs$ are not sufficient to evaluate the background contamination in the ϕ_{plane} distribution after the charge identification with 250 fb^{-1} . In Fig. 7, we can see the weak linear correlation between a and \tilde{b} . The mirror symme-

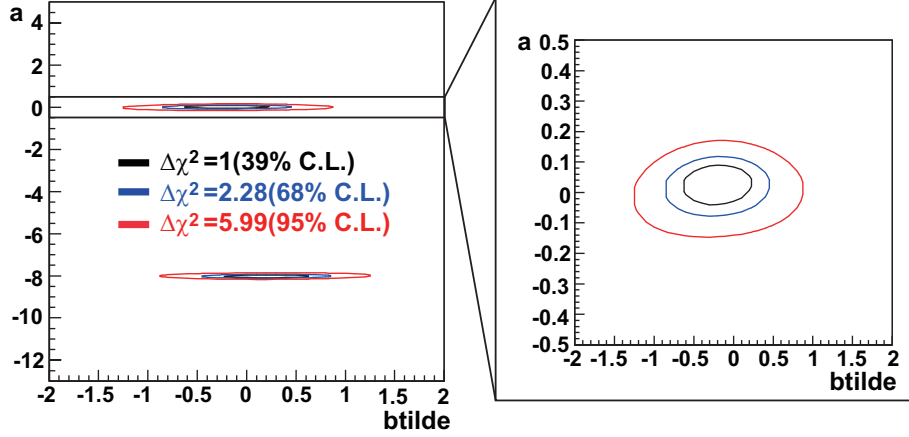


Figure 7: Contours similar to Fig.6 plotted in the b - \tilde{b} plane. Here, the integrated luminosity is assumed as $1,000 \text{ fb}^{-1}$.

try corresponding to $\tilde{b} \rightarrow -\tilde{b}$ is thus broken, although we still have the rotational symmetry corresponding to the combined transformation $(a, \tilde{b}) \rightarrow (-8.2 - a, -\tilde{b})$.

4.2 Theoretical consideration

Here we discuss theoretical possibilities to induce the non-renormalizable interactions in Eq. (2.1) as low energy effective theory. A simple setup is the Randall-Sundrum model [24], in which the gauge hierarchy problem can be solved by virtue of the warped metric. In the effective four dimensional theory of this model, the effective cutoff at the TeV scale emerges as $\Lambda = M_P \omega$ with the reduced Planck scale $M_P = 2.4 \times 10^{18} \text{ GeV}$ and the so-called warp factor $\omega \sim 10^{-15}$ associated with the warped metric. In this model, we can introduce effective higher-dimensional interactions [2]

$$\mathcal{L}_{\text{int}} = \sum_{A=1,2} c_A \frac{\phi^\dagger \phi}{\Lambda^2} \text{tr} [\mathcal{F}_A^{\mu\nu} \mathcal{F}_{A\mu\nu}], \quad (4.6)$$

where ϕ is the SM Higgs doublet field, the constants c_A are dimensionless parameters, and \mathcal{F}_A are the field strengths of the corresponding SM gauge groups, $U(1)_Y$ and $SU(2)_L$. After EW symmetry breaking, Eq. (4.6) is rewritten as interactions of the Higgs boson with photons, Z - and W -bosons,

$$\begin{aligned} \mathcal{L}_{\text{int}} = & \frac{c_{WW}}{\Lambda} \left(\frac{v}{\Lambda} \right) H W_{\mu\nu}^+ W^{-\mu\nu} + \frac{c_{ZZ}}{2\Lambda} H Z^{\mu\nu} Z_{\mu\nu} \\ & + \frac{c_{\gamma\gamma}}{\Lambda} \left(\frac{v}{\Lambda} \right) F^{\mu\nu} F_{\mu\nu} + \frac{c_{Z\gamma}}{2\Lambda} \left(\frac{v}{\Lambda} \right) Z^{\mu\nu} F_{\mu\nu}, \end{aligned} \quad (4.7)$$

where $W_{\mu\nu}^+$, $Z^{\mu\nu}$ and $F^{\mu\nu}$ are the field strengths of the W -boson, Z -boson and photon respectively. The couplings c_{WW} etc. can be described in terms of the original two

couplings c_1 , c_2 and the weak mixing angle θ_w as follows:

$$\begin{aligned}
c_{WW} &= c_2, \\
c_{ZZ} &= c_1 \sin^2 \theta_w + c_2 \cos^2 \theta_w, \\
c_{\gamma\gamma} &= c_1 \cos^2 \theta_w + c_2 \sin^2 \theta_w, \\
c_{Z\gamma} &= (-c_1 + c_2) \sin \theta_w \cos \theta_w.
\end{aligned} \tag{4.8}$$

We can identify $b = c_{WW}v/\Lambda$. In the same way, we can obtain \tilde{b} in Eq. (2.1) by effective interactions

$$\mathcal{L}_{\text{int}} = \sum_A \tilde{c}_A \frac{\phi^\dagger \phi}{\Lambda^2} \epsilon^{\mu\nu\rho\sigma} \text{tr} [\mathcal{F}_{A\mu\nu} \mathcal{F}_{A\rho\sigma}]. \tag{4.9}$$

We can also introduce

$$\mathcal{L}_{\text{int}} = c_0 \frac{\phi^\dagger \phi}{\Lambda^2} (D^\mu \phi)^\dagger (D_\mu \phi), \tag{4.10}$$

which gives rise to $a = c_0 v/\Lambda$ in Eq. (2.1). If we allow $c_A, \tilde{c}_A, c_0 = \pm \mathcal{O}(10)$, the parameters, a, b, \tilde{b} can be as large as order unity for $\Lambda = 1$ TeV.

5 Summary

In this work we have studied the sensitivity of the ILC to dimension-5 anomalous Higgs boson couplings to W^\pm pairs, using the decay $H \rightarrow WW^*$. We assumed a Higgs boson mass of 120 GeV, so that the Higgs boson will be produced in association with a Z boson through the Higgs-strahlung process. In order to avoid additional combinatorial backgrounds, we required the associated Z boson to decay invisibly into $\nu\bar{\nu}$ pairs. A direct measurement of this type is not expected to be possible at the LHC, due to the large QCD background.

Around the SM point $a = b = \tilde{b} = 0$, the coefficients of the anomalous couplings can be measured directly at the ILC by examining kinematic distributions such as the on-shell W boson momentum and the angle between W boson decay planes. Such a measurement will be able to probe the Lorentz structure of the HW^+W^- vertex, providing a direct test of the mechanism of spontaneous symmetry breaking. The sensitivity may be enhanced by combining the $ZH \rightarrow \nu\bar{\nu}WW^*$ channel considered here with other decay modes of Z such as $Z \rightarrow \ell\ell$ and $Z \rightarrow q\bar{q}$, although the large combinatorial background is expected to limit the sensitivity in the case of $Z \rightarrow q\bar{q}$.

Although we typically consider the anomalous couplings to be small corrections to the SM term, in principle a SM-like measurement is compatible with two distinct regions around $a = 0$ and $a \simeq -8.2$. This corresponds to a sign change in the coupling constant of the SM HW^+W^- vertex, which is only observable once a sign convention has been fixed by some other coupling (for example, the Yukawa coupling to b -quarks). The sign of the SM term could in principle be measured in interference effects.

Acknowledgments

The authors would like to thank all the members of the ILD detector optimization working group [25] and ILC physics subgroup [26] for useful discussions. This work is supported in part by the Creative Scientific Research Grant (No. 18GS0202) of the Japan Society for Promotion of Science, JSPS Grant-in-Aid for Scientific Research (No. 22244031) and the DOE Grants (No. DE-FG02-10ER41714). RNH is grateful to the Japan Society for the Promotion of Science for support during the initial stages of this project.

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